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JET CROSS SECTIONS IN POLARIZED PHOTON–HADRON COLLISIONS

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Abstract

We present a computation of one- and two-jet cross sections in polarized photon–hadron collisions, which is accurate to next-to-leading order in QCD. Our results can be used to compute photoproduction cross sections in electron–proton scattering. To this purpose, we investigate the structure of the polarized Weizsäcker–Williams function, where we include a universal, non-logarithmic term, neglected in the literature. We construct a Monte Carlo code, within the framework of the subtraction method, and we use it to study the phenomenology of jet production in the energy range relevant to HERA. In particular, we investigate the perturbative stability of our results, and we discuss the possibility of constraining polarized parton densities of the proton and the photon using jet data.

The last decade has seen an important advance in our understanding of polarized nucleon structure functions as a result of the analysis of deep inelastic scattering (DIS) data [1]. Unfortunately, the use of DIS data alone does not allow an accurate determination of the polarized parton densities. This is true in particular for the gluon, since this quantity contributes to DIS in leading order (LO) only via the Q^2 -dependence of the spin asymmetry (A_1^N), which could not be thoroughly studied experimentally so far. In the near future, the improved capabilities of handling polarized beams will allow high-energy colliders to operate in polarized mode, thus giving the possibility to study polarized scattering in a way complementary to DIS. At the RHIC collider at BNL, the first proton–proton polarized collisions are expected in two years from now. The possibility that the electron–proton HERA collider at DESY will run with polarized beams has also been considered for a long time [2].

At variance with DIS, collider physics offers a relatively large number of processes whose dependence upon the gluon density is dominant already at LO. The study of these processes is therefore crucial in order to measure this density in a direct way. Among the various high transverse momentum reactions with a sizeable gluonic contribution, jet production is an obvious candidate, because of the large rates. At HERA, the largest cross sections will be obtained by retaining those events where the electron is scattered at very small angles with respect to the beam line. In this way, the photon exchanged between the electron and the proton is almost on-shell, and the underlying dynamics is that of a photoproduction process.

As was shown in ref. [3], a polarized version of the HERA collider with $\sqrt{S} \approx 300$ GeV would be a very promising and useful facility for studying polarized photoproduction reactions. In particular, jet production cross sections display a strong sensitivity to the polarized gluon distribution of the proton. However, the analysis of ref. [3] is based on a lowest order QCD computation. From unpolarized physics it is well known that, in order to describe the data in a satisfactory way, the calculation of next-to-leading order (NLO) QCD corrections to the jet cross sections is mandatory.

The computation of the NLO QCD corrections to one- and two-jet production in polarized photon–hadron collisions is the purpose of this paper. As is well known, in QCD photoproduction cross sections are written as the sum of two terms (point-like and hadronic components), neither of which is physically meaningful by itself. The hadronic part, in which the photon behaves like a hadron, can be accounted for by using the available results for jet production in polarized hadron–hadron scattering [4]. The NLO QCD corrections for the point-like part are presented in this paper for the first time. Such a calculation needs the one-loop $2 \rightarrow 2$ and tree-level $2 \rightarrow 3$ polarized amplitudes as input (one of the particles in the initial state being the photon and the other a quark or gluon). These amplitudes are already known [5]. The next step is to implement them in a Monte Carlo code which allows the calculation of any three-parton infrared-safe observable. In order to do this, we will use the general algorithm, based on the subtraction method, presented in ref. [6]. Following the same strategy as already adopted in the case of polarized hadron–hadron collisions [4], we build a modified version of the Monte Carlo code of ref. [7], which deals with unpolarized photon–hadron collisions. The reader can find further details in ref. [4].

In photoproduction processes initiated by an electron, the electron can be considered to be equivalent to a beam of real photons, whose distribution in energy (Weizsäcker–Williams

function [8]) can be computed theoretically. The polarized electron–hadron cross section reads

$$d\Delta\sigma_{eH}(K_e, K_H) = \int_{y_{min}}^{y_{max}} dy \Delta f_\gamma^{(e)}(y) d\Delta\sigma_{\gamma H}(yK_e, K_H), \quad (1)$$

where K_e and K_H are the momenta of the incoming electron and hadron respectively, and $0 < y_{min} < y_{max} \leq 1$ are fixed by kinematical boundary conditions or experimental cuts. The quantities $\Delta\sigma$ appearing in eq. (1) are defined in terms of cross sections $\sigma(\lambda_1, \lambda_2)$ for incoming particles of definite helicities

$$d\Delta\sigma = \frac{1}{4} \left(d\sigma(+, +) + d\sigma(-, -) - d\sigma(+, -) - d\sigma(-, +) \right). \quad (2)$$

The unpolarized electron–hadron cross section can be obtained from eq. (1) with the formal substitutions

$$\Delta\sigma \longrightarrow \sigma, \quad \Delta f_\gamma^{(e)} \longrightarrow f_\gamma^{(e)}. \quad (3)$$

As mentioned before, the polarized and unpolarized Weizsäcker–Williams functions ($\Delta f_\gamma^{(e)}$ and $f_\gamma^{(e)}$, respectively) can be computed. As far as the unpolarized case is concerned, we use the form [9]

$$f_\gamma^{(e)}(y) = \frac{\alpha_{em}}{2\pi} \left[\frac{1 + (1-y)^2}{y} \log \frac{Q^2(1-y)}{m_e^2 y^2} + 2m_e^2 y \left(\frac{1}{Q^2} - \frac{1-y}{m_e^2 y^2} \right) \right]. \quad (4)$$

Notice that eq. (4) contains a non-logarithmic term with a singular behaviour for $y \rightarrow 0$. In ref. [10] this term was shown to give non-negligible contributions (of about 7% and 5% for $Q^2 = 0.01 \text{ GeV}^2$ and $Q^2 = 4 \text{ GeV}^2$, respectively) in the case of unpolarized jet production at HERA. One expects an analogous non-logarithmic term to appear also in the polarized Weizsäcker–Williams function. To the best of our knowledge, however, such a term has never been presented in the literature. We therefore computed $\Delta f_\gamma^{(e)}$ from scratch. We closely follow the procedure of ref. [9], in which the interested reader will find further details. We consider the generic production process

$$e(p; n) + a(k; s) \longrightarrow e(p') + X, \quad q = p - p', \quad (5)$$

where a is a parton whose momentum k , which is kept off-shell for the time being, will eventually be put on-shell, and n, s are spin vectors. In the case of polarized scattering, we only need the antisymmetric part of the partonic and leptonic tensors (see for example ref. [11]):

$$W_A^{\mu\nu} = \frac{W^{\mu\nu} - W^{\nu\mu}}{2} = i \frac{\sqrt{|k^2|}}{k \cdot q} \epsilon^{\mu\nu\rho\sigma} q_\rho \left[g_1(q^2, k \cdot q) s_\sigma + g_2(q^2, k \cdot q) \left(s_\sigma - \frac{q \cdot s}{q \cdot k} k_\sigma \right) \right], \quad (6)$$

$$T_A^{\mu\nu} = \frac{T^{\mu\nu} - T^{\nu\mu}}{2} = -2im_e \epsilon^{\mu\nu\rho\sigma} n_\rho q_\sigma, \quad (7)$$

where g_1 and g_2 are the usual structure functions relevant to polarized collisions. The spin vectors for the incoming electron and parton can be written as

$$n_\mu = N_n \left(p_\mu - \frac{m_e^2}{k \cdot p} k_\mu \right), \quad s_\mu = N_s \left(k_\mu - \frac{k^2}{k \cdot p} p_\mu \right), \quad (8)$$

where N_n and N_s are normalization factors such that $|n^2| = |s^2| = 1$. The spin vectors satisfy the transversality conditions $n \cdot p = 0$ and $s \cdot k = 0$. The electron-parton cross section follows from direct computation¹ ($y = (k \cdot q)/(k \cdot p)$, $k^2 \rightarrow 0$):

$$\begin{aligned} d\Delta\sigma_{ea}(p, k) &= \frac{1}{4k \cdot p} \frac{e^2 W_A^{\mu\nu} T_{A\mu\nu}}{q^4} \frac{d^3 p'}{(2\pi)^2 2E'} \\ &= -\frac{\alpha_{\text{em}}}{2\pi} \left[\frac{1 - (1-y)^2}{yq^2} + \frac{2m_e^2 y^2}{q^4} \right] \frac{g_1(0, k \cdot q)}{4y k \cdot p} dq^2 dy. \end{aligned} \quad (9)$$

Notice that the terms proportional to g_2 drop from this equation. This expression can be related to the polarized cross section for real photon-parton scattering, which can be computed by convoluting the parton tensor of eq. (6) with the antisymmetric part of the photon polarization density matrix,

$$P_A^{\mu\nu} = \frac{1}{2} (\varepsilon^\mu \varepsilon^{\nu*} - \varepsilon^\nu \varepsilon^{\mu*}) = \frac{i}{2\sqrt{|q^2|}} \epsilon^{\mu\nu\rho\sigma} q_\rho t_\sigma, \quad (10)$$

where ε^μ is the photon polarization vector and t^μ is its spin vector

$$t_\mu = N_t \left(q_\mu - \frac{q^2}{k \cdot q} k_\mu \right). \quad (11)$$

We get ($q^2 \rightarrow 0$)

$$\Delta\sigma_{\gamma a}(q, k) = \frac{1}{4k \cdot q} W_A^{\mu\nu} P_{A\mu\nu} = \frac{g_1(0, k \cdot q)}{4k \cdot q}. \quad (12)$$

Combining eqs. (9) and (12), and integrating over q^2 in the accessible kinematical range (see ref. [9]), we get

$$d\Delta\sigma_{ea}(p, k) = \Delta\sigma_{\gamma a}(yp, k) \Delta f_\gamma^{(e)}(y) dy, \quad (13)$$

where

$$\Delta f_\gamma^{(e)}(y) = \frac{\alpha_{\text{em}}}{2\pi} \left[\frac{1 - (1-y)^2}{y} \log \frac{Q^2(1-y)}{m_e^2 y^2} + 2m_e^2 y^2 \left(\frac{1}{Q^2} - \frac{1-y}{m_e^2 y^2} \right) \right]. \quad (14)$$

We therefore find that also in the case of polarized scattering a non-logarithmic term is present. This term, at variance with the unpolarized case, is not singular for $y \rightarrow 0$. However, exactly as in the unpolarized case, its behaviour for small y is equal to the behaviour of the term that multiplies the logarithm.

We can now go back to eq. (1), and compute $d\Delta\sigma_{\gamma H}$. Using the factorization theorems of QCD, we can write

$$d\Delta\sigma_{\gamma H}(K_\gamma, K_H) = d\Delta\sigma_{\gamma H}^{\text{point}}(K_\gamma, K_H) + d\Delta\sigma_{\gamma H}^{\text{hadr}}(K_\gamma, K_H), \quad (15)$$

where

$$d\Delta\sigma_{\gamma H}^{\text{point}}(K_\gamma, K_H) = \sum_j \int dx \Delta f_j^{(H)}(x, \mu'_H) d\Delta\hat{\sigma}_{\gamma j}(K_\gamma, xK_H, \alpha_s(\mu'_R), \mu'_R, \mu'_H, \mu_\gamma), \quad (16)$$

$$\begin{aligned} d\Delta\sigma_{\gamma H}^{\text{hadr}}(K_\gamma, K_H) &= \sum_{ij} \int dx dy \Delta f_i^{(\gamma)}(x, \mu_\gamma) \Delta f_j^{(H)}(y, \mu''_H) \\ &\quad \times d\Delta\hat{\sigma}_{ij}(xK_\gamma, yK_H, \alpha_s(\mu''_R), \mu''_R, \mu''_H, \mu_\gamma). \end{aligned} \quad (17)$$

¹We keep only the first term in the expansion of g_1 and g_2 in series of q^2 : the remaining terms give non-factorizable corrections which are suppressed by powers of Q^2/M_X^2 .

We remind the reader that the two terms in the RHS of eq. (15) (denoted as point-like and hadronic components) are not separately well defined in perturbative QCD beyond leading order: only their sum is physically meaningful. The polarized parton distribution functions $\Delta f_i^{(H)}$ and $\Delta f_i^{(\gamma)}$ are defined, as usual, in terms of densities for partons of definite helicity in hadrons of definite helicity ($A = H, \gamma$):

$$\Delta f_i^{(A)} = f_{i+}^{(A+)} - f_{i-}^{(A+)} = f_{i-}^{(A-)} - f_{i+}^{(A-)}. \quad (18)$$

The subtracted partonic cross sections $d\Delta\hat{\sigma}_{\gamma j}$ and $d\Delta\hat{\sigma}_{ij}$ are finite at any order in QCD and, for the purpose of this paper, have been evaluated to NLO ($\alpha_{\text{em}}\alpha_s^2$ and α_s^3 , respectively). Notice that while the dependence upon the scales μ'_R, μ'_H and μ''_R, μ''_H cancels (up to NNLO terms) in the point-like and hadronic components separately, the dependence upon μ_γ in the point-like component is compensated by a corresponding dependence in the hadronic component. We finally point out that equations completely similar to eqs. (15)–(17) also hold for unpolarized scattering.

In what follows, we will present phenomenological predictions relevant to HERA physics. To this purpose, we will evaluate both the polarized and unpolarized cross sections. Since one of the major goals of high-energy colliders in the polarized mode will be that of studying the polarized parton densities, we will make a definite choice for the unpolarized densities throughout the paper, and study the dependence of our predictions for the asymmetries upon the polarized densities. Namely, we will adopt the set MRSA' [12] for the unpolarized proton, and the set GRV-HO [13] for the unpolarized photon. The value of Λ_{QCD} associated to MRSA' ($\Lambda_5^{\overline{\text{MS}}} = 152$ MeV) is well below the current world average; however, it is closer to the value of Λ_{QCD} associated to GRV-HO and to all the polarized sets we are going to use. The use of more modern proton sets would only affect the absolute normalization of the cross sections, without any sizeable change in the shape of the distributions.

As far as the polarized densities in the proton are concerned, DIS data available at present leave the gluon density largely unconstrained, while they determine the quark densities to a reasonable extent. Several parametrizations are therefore available, whose gluon densities are very different from each other. We will consider the following sets: GRSV STD and GRSV MAXG [14], DSS1, DSS2 and DSS3 [15], and GS-C [16] (see ref. [4] for a discussion about this choice and a comparison between the various sets). GRSV STD will be the default set when studying the perturbative stability of our results.

The polarized densities in the photon are completely unmeasured so far, and models for them have to be invoked. To obtain a realistic estimate for the theoretical uncertainties due to these densities, we use the two very different scenarios considered in ref. [17], assuming “maximal” ($\Delta f^{(\gamma)}(x, \mu_I^2) = f^{(\gamma)}(x, \mu_I^2)$) or “minimal” ($\Delta f^{(\gamma)}(x, \mu_I^2) = 0$) saturation of the positivity constraint² $|\Delta f^{(\gamma)}| \leq f^{(\gamma)}$ at the input scale μ_I ($\mathcal{O}(0.5$ GeV)) for the QCD evolution (here, the unpolarized densities are those of the set GRV-HO). The corresponding sets are denoted as SV MAX γ and SV MIN γ respectively. The former will be our default choice when studying perturbative stability.

In order to compute the NLO results consistently, we will use the NLO-evolved version of the previously mentioned parton distributions³. We point out that these NLO-evolved

²Strictly speaking, this bound is violated in QCD, and the deviations may become sizeable at small scales. For a discussion on this point, see for example ref. [18]

³In fact, for most of the sets considered here, there exists a LO-evolved version as well.

sets will also be used to compute the Born results; indeed, as observed in ref. [4], the use of LO-evolved parton distributions could introduce differences between the LO and NLO cross sections that are much larger than those that can be expected on the basis of perturbative considerations.

Since this paper presents the first NLO calculation of jet cross sections in polarized photon–hadron collisions, it is mandatory to assess the effect of the radiative corrections. Lacking an NNLO computation, the most reliable estimate of the theoretical uncertainty affecting our results comes from the study of the dependence upon the mass scales entering eqs. (16) and (17), at fixed values of the other input parameters. Unless otherwise specified, we will set all the scales equal to a common value μ , and vary μ in the range $\mu_0/2 \leq \mu \leq 2\mu_0$. Here μ_0 is the default scale, which we set equal to half of the total transverse energy of the event; other sensible choices for μ_0 give completely similar results. From eq. (1), it is clear that all the dynamical information on the ep process in the Weizsäcker–Williams approximation are contained in the γp cross section. Therefore, we will study the scale dependence of our results using monochromatic photon–proton collisions at $E_{cm}^{(\gamma p)} = 268$ GeV; the scale dependence at lower centre-of-mass energies, and in electron–proton collisions, is comparable or milder.

We start by considering the single-inclusive jet pseudorapidity, requiring $p_T > 10$ GeV, where p_T is the transverse momentum of the jet. The jet has been defined using the cone algorithm with $R = 1$. The results for polarized and unpolarized cross sections are presented in fig. 1. The size of the radiative corrections is larger in the case of unpolarized collisions; in the case of polarized collisions, the contributions from the various partonic subprocesses do not have the same sign, and cancellations occur. In both cases, the variation of the scale induces a variation of the cross section of the order of 10% over most of the η range considered. The scale dependence is strongly reduced when going from Born (boxes and circles) to NLO (dashed and dotted histograms) results. However, this is not true in the region of negative η 's, where the scale dependence of the Born cross section is very small, and smaller than that of the NLO results. This could be the signal of a failure of the perturbative expansion, which would be extremely relevant from the phenomenological point of view, since it is precisely from this region that one hopes to extract information on the gluon density in the proton [3]. In order to understand this point, we have studied the scale dependence of our results more carefully, by varying separately all the scales that enter eqs. (16) and (17). At the Born level, the hadronic component has a relatively mild μ_H'' dependence, a sizeable μ_γ dependence, and a very large μ_R'' dependence. Although the μ_γ and μ_R'' variations tend to compensate each other, the latter is dominant, and the overall result is a large scale dependence, which is precisely the effect we see in the positive η region in fig. 1, where the hadronic component is dominant. On the other hand, the point-like component displays a dependence upon μ_H' and μ_R' of about the same size, but opposite in sign (at the Born level, the point-like component does not depend upon μ_γ), and a simultaneous variation of the two scales leads to an almost perfect cancellation of the effects induced by them. This is the origin of the small scale dependence of the Born result that we see in the negative η region in fig. 1. When going to NLO, in the hadronic contribution the μ_H'' dependence basically reduces to zero, and the μ_R'' dependence is sizeably reduced; this is what we expect, as remarked after eq. (18). The μ_γ dependence is also reduced, although it remains sizeable; indeed, the effects of the variation of μ_γ are only partially cancelled in the hadronic component, since the complete cancellation (up to NNLO terms) takes place when the point-like component is included. In fact, we can observe this

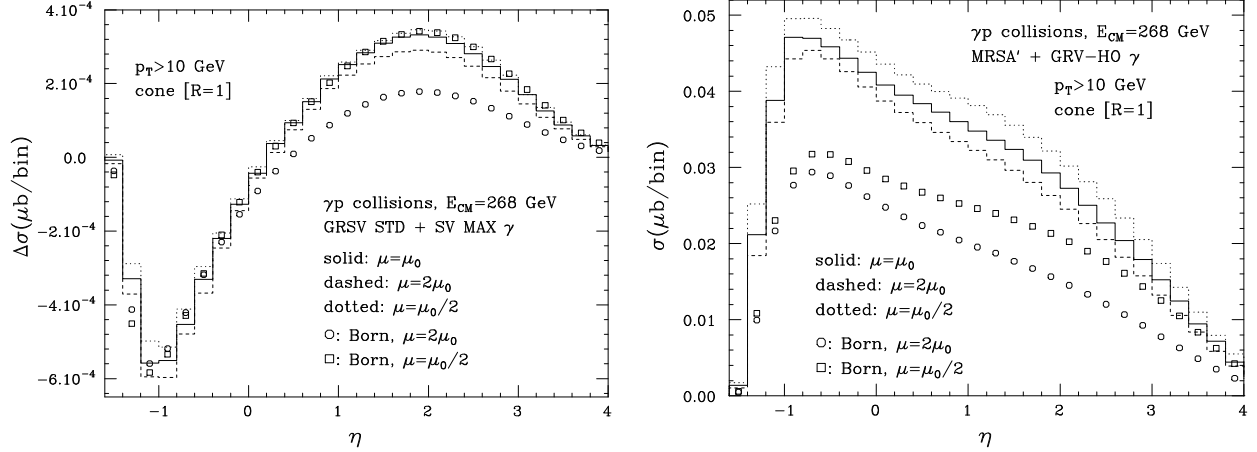


Figure 1: *Scale dependence of the single-inclusive jet pseudorapidity in polarized (left) and unpolarized (right) collisions.*

cancellation by looking at the μ_γ dependence of the point-like component at NLO, whose size is almost the same as that of the hadronic component, but opposite in sign. Finally, the μ'_H and μ'_R dependences of the point-like component are largely reduced with respect to the corresponding dependences at LO. In conclusion, the scale dependence of the η distribution is consistent with what we expect from perturbative considerations. The small scale dependence of the Born result in the negative η region is only due to an incidental cancellation, which takes place when fixing all the scales to the same value. A more sophisticated treatment of the scale dependence would result in a scale dependence larger at the Born level than at the NLO level.

We also studied several double differential observables, by considering the two leading jets of each event. As an example, we present in fig. 2 the invariant mass of the pair. We require the hardest (next-to-hardest) jet to have transverse momentum larger than 15 GeV (10 GeV). We do not impose any η cuts, in order to maximize the scale dependence (as shown in fig. 1, the scale dependence is larger for the smallest and largest accessible η values, since they correspond to small transverse momenta). Therefore, in the case of η cuts simulating a realistic geometrical acceptance, the scale dependence of our results would be smaller. We see that the case of the invariant mass is pretty similar to the case of single-inclusive pseudorapidity. The size of radiative corrections is larger in the case of unpolarized scattering, and the scale dependence of the polarized result is comparable to that of the unpolarized one. A reduction of the scale dependence can be observed when going from LO to NLO, except for the region close to the threshold, where there is no contribution at the Born level (due to the asymmetric transverse-momentum cuts), and the NLO result is effectively a LO one.

For all the single-inclusive and double-differential jet observables that we have studied, the same pattern as outlined above is reproduced. The size of radiative corrections is usually smaller in the polarized case than in the unpolarized case. The inclusion of NLO terms reduces the magnitude of the scale dependence with respect to the LO results, in both the polarized and unpolarized cases. When considering double-differential cross sections, special care has

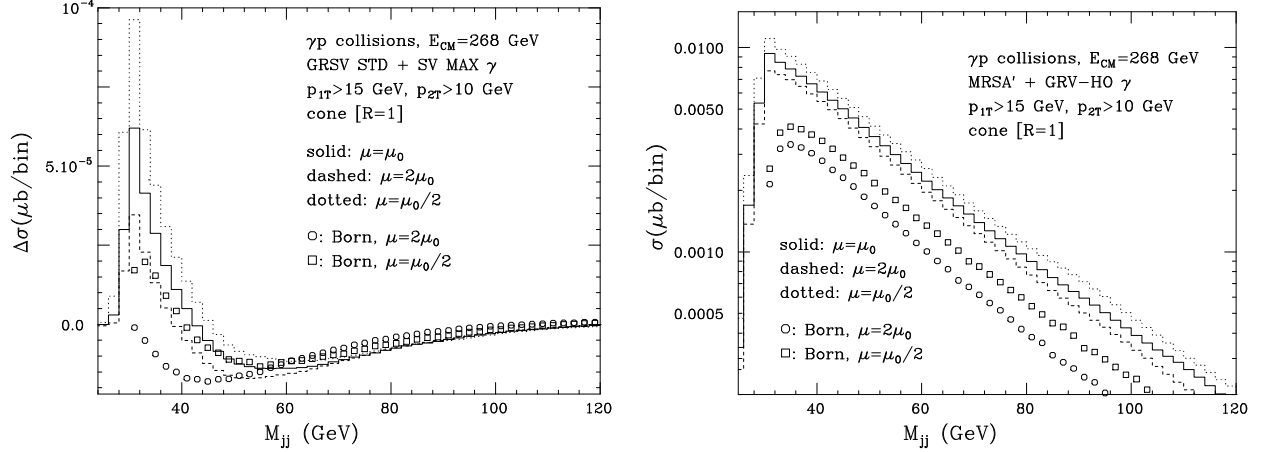


Figure 2: *Scale dependence of the invariant mass of the jet-jet pair in polarized (left) and unpolarized (right) collisions.*

to be taken in the case where the cuts on the transverse momenta of the two leading jets are set to the same value, since there are regions in the phase space where the perturbative computation fails to give a sensible prediction. This situation has been discussed at length in ref. [10] for unpolarized jet production. We did not find any major difference in the case of polarized scattering. We also investigated the effect of changing the jet definition, by considering the k_T algorithm proposed in ref. [19]. This basically only amounts to a change in the normalization, while the effects due to the variation of the scales are of the same size as those studied above. Finally, the perturbative stability of the results obtained with different parton densities is identical to the one presented here.

We now turn to the problem of studying the dependence of our results upon the available proton and photon polarized parton densities. To this end, we will consider electron-proton collisions in the Weizsäcker-Williams approximation, setting $E_{cm}^{(ep)} = 300$ GeV, $y_{min} = 0.2$, $y_{max} = 0.8$, and $Q^2 = 0.01$ GeV² (see eqs. (1), (4) and (14)). All the scales will be set to the default value. In fig. 3 we present the results for the asymmetry

$$\mathcal{A}_\eta = \frac{d\Delta\sigma_{ep}/d\eta}{d\sigma_{ep}/d\eta}, \quad (19)$$

where η is the pseudorapidity of the single-inclusive jet. A cut of $p_T > 10$ GeV has been applied. On the LHS of the figure, we have chosen GRSV STD as the polarized proton set, and we evaluated the results for both the polarized photon sets we consider in this paper. The Born and NLO results are both shown. We can see that in the large (positive) η region the difference induced by the choice of the two photon sets is extremely large. On the other hand, towards negative η values this difference tends to vanish. This is because in that region the point-like component, which does not depend upon photon densities, is the dominant one. We can also observe that in the positive η region there is a very small difference between the NLO and LO results, while for negative η 's the radiative corrections are positive and reduce the asymmetry considerably. It is worth noticing that, as can be seen from fig. 1, this reduction

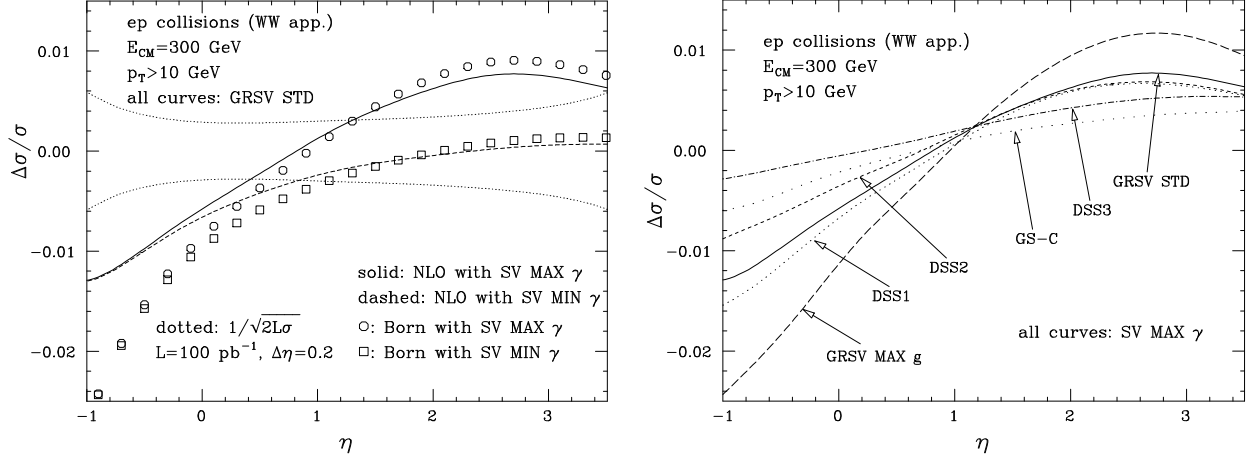


Figure 3: *Asymmetries versus pseudorapidity in single-inclusive jet production, for electron-proton collisions at $E_{cm} = 300$ GeV. The predictions for several photon (left) and proton (right) densities are shown.*

is mainly due to the increase of the unpolarized cross section at NLO with respect to LO, while the polarized result is basically unchanged. On the RHS of fig. 3, we show the curves obtained by fixing the polarized photon set to SV MAX γ , and by considering the various polarized proton sets. As expected, the largest differences can be seen at negative η values, where theoretical predictions can vary for about one order of magnitude. The shape of the asymmetries is quite similar for all the sets, except for GS-C; this reflects the fact that the shape of the GS-C gluon is rather different from the shape of the gluons of the other sets. The same effect has indeed been observed in polarized pp collisions [4].

On the LHS of fig. 3 we also show (dotted lines) an estimate of the minimum value of the asymmetry observable at HERA. This is calculated using the following formula

$$(\mathcal{A}_\eta)_{min} = \frac{1}{P_e P_p} \frac{1}{\sqrt{2\sigma_{ep}\mathcal{L}\epsilon}}, \quad (20)$$

where \mathcal{L} is the integrated luminosity, $P_{e(p)}$ is the polarization of the electron (proton) beam, and the factor $\epsilon \leq 1$ accounts for experimental efficiencies; σ_{ep} is the unpolarized cross section integrated over a range in rapidity ($\Delta\eta$), which has been chosen, for the present case, equal to 0.2. In fig. 3 we used $\mathcal{L} = 100 \text{ pb}^{-1}$, $P_e = P_p = 1$ and $\epsilon = 1$. It follows that, if high luminosity will be collected, it will be possible to get information on the polarized parton densities in the proton, given the fact that realistic values for the polarization of the beams will be such that $P_e P_p \simeq 0.5$. Notice that the cross section corresponding to sets such as DSS3 or GS-C will be hardly measurable at HERA, regardless of the luminosity collected. Of course, the bin size $\Delta\eta$ can be enlarged; this would result in a smaller observable asymmetry, at the price of a loss in resolution. As far as the polarized photon densities are concerned, if the “real” densities are similar to those of the set SV MIN γ , it will be extremely hard to even get the experimental evidence of a hadronic contribution to the polarized cross section. On the other hand, a set like SV MAX γ appears to give measurable cross sections, but this

conclusion strongly depends upon the polarized proton densities. As observed in ref. [3], sensible information on the polarized photon densities will be obtained only after a more precise knowledge of the proton densities will be available. At this point, we have to comment upon the effect of the non-logarithmic term in the polarized Weizsäcker–Williams function. It turns out that, for the single-inclusive distribution we are studying here, this term reduces the absolute value of the cross section by a factor of about 6%. Since the non-logarithmic term in the unpolarized case also reduces the cross section by a factor of about 7% [10], it follows that the asymmetry is to a good extent independent from the presence of non-logarithmic terms in the Weizsäcker–Williams functions. Finally, if a different jet definition is adopted, the asymmetries change for less than 1%; this feature has already been observed in the case of pp collisions [4].

In principle, it is conceivable to study asymmetries also for monochromatic photon–proton collisions. From the experimental point of view, this would correspond to selecting events with the Weizsäcker–Williams y (that is, the fraction of electron energy carried away by the photon) in a given narrow range (which corresponds to taking y_{min} and y_{max} very close to each other in eq. (1)). If we compute the asymmetry of eq. (19) with $(\Delta)\sigma_{\gamma p}$ instead of $(\Delta)\sigma_{ep}$, the values we obtain are larger than those shown in fig. 3. For $E_{cm}^{(\gamma p)} = 134$ GeV, the asymmetry can be twice as large as that obtained using eq. (19). Unfortunately, this is not sufficient. Indeed, what is actually measured at HERA is an electron–proton cross section. Therefore, to get a physically observable quantity one should multiply the γp asymmetry by the ratio of the integral of the polarized Weizsäcker–Williams function in the range $y_{min} \leq y \leq y_{max}$ over the same integral calculated for the unpolarized function. This ratio tends to zero for $y_{min}, y_{max} \rightarrow 0$ ($y_{max} - y_{min}$ is kept fixed), thus disfavours the measurement of the asymmetry at small photon–proton centre-of-mass energies.

In conclusion, we reported the first calculation of jet cross sections in polarized photon–hadron collisions, which is accurate to NLO in perturbative QCD. For all the observables considered, it has been found that the size of the radiative corrections is moderate, and that the scale dependence is smaller than that of the LO result. The inclusion of the NLO terms reduces the size of the asymmetries in the pseudorapidity region where the contamination from the hadronic photon contribution is minimal. In order to get information on the polarized gluon densities in the proton and in the photon, a very high luminosity must be collected at HERA. We also calculated the universal, non-logarithmic term in the polarized Weizsäcker–Williams function. Its contribution to the cross sections is small with respect to the logarithmic term, analogously to what happens in the unpolarized case.

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References

- [1] For a compilation of references to the data, see P.J. Mulders and T. Sloan, Summary talk of Spin Physics Working Group at the 6th International Workshop on “Deep Inelastic Scattering and QCD”, Brussels (1998), [hep-ph/9806314](#).
- [2] See e.g. Proc. of the 1997 Workshop on “Physics with Polarized Protons at HERA”, eds. A. De Roeck and T. Gehrmann, Hamburg, Germany, DESY-Proceedings-1998-01.
- [3] M. Stratmann and W. Vogelsang, *Z. Phys.* **C74** (1997) 641; Proc. of the 1995/96 Workshop on “Future Physics at HERA”, Hamburg, Germany, eds. G. Ingelman, A. de Roeck, and R. Klanner, p. 815;
J.M. Butterworth, N. Goodman, M. Stratmann and W. Vogelsang, Proc. of the 1997 Workshop on “Physics with Polarized Protons at HERA”, Hamburg, Germany, CERN-TH/97-310, [hep-ph/9711250](#).
- [4] D. de Florian, S. Frixione, A. Signer and W. Vogelsang, *Nucl. Phys.* **B539** (1999) 455.
- [5] D. de Florian and W. Vogelsang, *Phys. Rev.* **D57** (1998) 4376.
- [6] S. Frixione, Z. Kunszt and A. Signer, *Nucl. Phys.* **B467** (1996) 399.
- [7] S. Frixione, *Nucl. Phys.* **B507** (1997) 295.
- [8] C.F. Weizsäcker, *Z. Phys.* **88** (1934) 612; E.J. Williams, *Phys. Rev.* **45** (1934) 729.
- [9] S. Frixione, M.L. Mangano, P. Nason and G. Ridolfi, *Phys. Lett.* **B319** (1993) 339.
- [10] S. Frixione and G. Ridolfi, *Nucl. Phys.* **B507** (1997) 315.
- [11] A. Gabrieli and G. Ridolfi, *Phys. Lett.* **B417** (1998) 369.
- [12] A. D. Martin, R. G. Roberts and W. J. Stirling, *Phys. Lett.* **B354** (1995) 155.
- [13] M. Glück, E. Reya and A. Vogt, *Phys. Rev.* **D46** (1992) 1973.
- [14] M. Glück, E. Reya, M. Stratmann and W. Vogelsang, *Phys. Rev.* **D53** (1996) 4775.
- [15] D. de Florian, O.A. Sampayo and R. Sassot, *Phys. Rev.* **D57** (1998) 5803.
- [16] T. Gehrmann and W.J. Stirling, *Phys. Rev.* **D53** (1996) 6100.
- [17] M. Stratmann and W. Vogelsang, *Phys. Lett.* **B386** (1996) 370.
- [18] S. Forte, G. Altarelli and G. Ridolfi, Proc. of “QCD 98”, Montpellier, July 1998, CERN-TH/98-285, [hep-ph/9808462](#).
- [19] S. D. Ellis and D. E. Soper, *Phys. Rev.* **D48** (1993) 3160.